

Leptogenesis and rescattering in supersymmetric models

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Abstract

The observed baryon asymmetry of the Universe can be due to the $B-L$ violating decay of heavy right handed (s)neutrinos. The amount of the asymmetry depends crucially on their number density. If the (s)neutrinos are generated thermally, in supersymmetric models there is limited parameter space leading to enough baryons. For this reason, several alternative mechanisms have been proposed. We discuss the nonperturbative production of sneutrino quanta by a direct coupling to the inflaton. This production dominates over the corresponding creation of neutrinos, and it can easily (i.e. even for a rather small inflaton-sneutrino coupling) lead to a sufficient baryon asymmetry. We then study the amplification of MSSM degrees of freedom, via their coupling to the sneutrinos, during the rescattering phase which follows the nonperturbative production. This process, which mainly influences the (MSSM) D -flat directions, is very efficient as long as the sneutrinos quanta are in the relativistic regime. The rapid amplification of the light degrees of freedom may potentially lead to a gravitino problem. We estimate the gravitino production by means of a perturbative calculation, discussing the regime in which we expect it to be reliable.

1 Introduction and summary

The generation of the Baryon Asymmetry of the Universe (BAU) [1] represents one of the puzzles of Cosmology. Three ingredients are required [2] to achieve this task: baryon number violation, C and CP violation, and departure from thermal equilibrium. The baryon number violation can be challenging to implement, because it must be consistent with the current lower bound on the proton lifetime, $\tau_p \gtrsim 10^{32}$ years [3]. The Standard Model (SM) is a C and CP violating theory, and contains non-perturbative $B+L$ violating interactions (sphalerons) [4]—which are rapid in the early Universe but unable to mediate proton decay. However, it seems difficult to use this baryon number violation to create the asymmetry in the SM [5] and its more popular extensions [6]. An

number (L)-violating out-of-equilibrium interaction, and then allow the sphalerons to reprocess part of it into a baryon asymmetry. An appealing feature of this scenario is that while neutrino masses are experimentally observed [8] (and are L -violating, if they are Majorana), there is still no evidence for baryon number violation.

The above idea is naturally implemented [7] in the context of the see-saw [9], which is a minimal mechanism for generating neutrino masses much smaller than the ones of the charged leptons. Three right-handed (r.h.) neutrinos N_i are added to the SM particle content, given Yukawa interactions with the lepton and Higgs doublets, and large Majorana masses. This gives the three light neutrinos very small masses, due to their small mixing with the heavy r.h. neutrinos through the Dirac mass. Grand Unified Theories (GUT) and their supersymmetric versions, that constitute natural candidates for the Physics beyond the SM, often contain r.h. neutrinos in their particle content. In this paper we consider the supersymmetric version of the see-saw mechanism, which is theoretically attractive because it addresses the hierarchy between the Higgs and r.h. neutrino masses.

The r.h. (s)neutrinos of the see-saw can generate the BAU via leptogenesis, in a three steps process [7, 10, 11, 12]. First, some (CP symmetric) number density of (s)neutrinos is created in the early Universe. Then, a lepton asymmetry is generated in their CP violating out-of-equilibrium decay. Finally, the lepton asymmetry is partially reprocessed into a baryon one by the $B + L$ violating interactions, provided it is not washed out by lepton number violating scatterings. In this paper, we are mostly interested in the first step, although in the next section we will also briefly review the decay and washout processes.

The most straightforward and cosmological model-independent mechanism to generate r.h. (s)neutrinos is via scattering in the thermal bath [12]. However, as discussed in section 2, the parameter space available is restricted in supergravity-motivated models. Indeed, unless some enhancement of the CP asymmetry characterizing the r.h. (s)neutrino decay is present (which occurs for example if they are nearly degenerate in mass) the generation of a sufficient lepton number poses a rather strong lower bound on their mass [13, 14, 15]. For thermal production, this translates into a lower bound on the reheating temperature T_{RH} of the thermal bath. On the other hand, if supergravity is assumed, T_{RH} cannot be taken arbitrarily large without leading to an overproduction of gravitinos [16, 17, 18, 19, 20, 21]. The two requirements are compatible only provided a nearly maximal CP asymmetry (again, banning any possible enhancement from mass degeneracy) is present in the r.h. (s)neutrino decay. If this is not the case, alternative mechanisms of production for the r.h. (s)neutrino have to be considered.

As remarked in [22], leptogenesis can be achieved if at least one r.h. sneutrino has a smaller mass than the Hubble parameter, *i.e.* ¹ $M_N < H$, during inflation. In this case, quantum fluctuations of this sneutrino component are produced during the inflationary expansion, and amplified to generate a classical condensate. The decay of the condensate eventually generates the required lepton asymmetry. The above requirement $M_N < H$ is not trivially satisfied in a supergravity context, since supergravity

¹We will use N as a shorthand for the superfield, its scalar and fermionic component. We explicitly refer to the particle type when this could cause a confusion.

model [23]. In this case, however, a suitable choice of the Kähler potential can induce a negative mass term $m_{\text{ind}}^2 \simeq -H^2$, so that a large expectation value will be generated for the sneutrino component during inflation [23]. This also leads to the formation of a condensate during inflation, and to successful leptogenesis as in the previous case.

Large variances can be produced during inflation if the sneutrinos are not too strongly coupled to the inflaton field ϕ , since this would generate a high effective mass which could fix $\langle \tilde{N} \rangle = 0$ during inflation. However, if one of the r.h. (s)neutrinos is coupled to the inflaton, there is the obvious possibility that a sufficient amount of (s)neutrino quanta is generated when the inflaton decays. Quite remarkably, for a rather wide range of models this decay occurs in a nonperturbative way [24, 25] (this is known in the literature as *preheating* [25]). In models of chaotic inflation [26], this is due to the coherent oscillations of the inflaton field, which can be responsible for a parametric amplification of the bosonic fields to which the inflaton is coupled². It is important to remark that this resonant amplification does not require very high couplings between the inflaton and the produced fields. For a coupling of the form $(g^2/2)\phi^2 N^2$ in the scalar potential, resonant amplification of the field N already occurs for $g^2 \gtrsim 10^{-8}$ [30], if the mass of N is negligible at the end of inflation, and if a massive inflaton is considered. For a massless inflaton ($V(\phi) = \lambda\phi^4/4$), an efficient resonance is present also for much smaller values of g (we will show this explicitly in section 3), since in this case the resonance is not halted by the expansion of the Universe [31].

If the produced particle is very massive ($M_N \gtrsim m_\phi$), the effectiveness of the resonance becomes a highly model dependent issue. A potential of the form $V(\phi) + M_N^2 N^2/2 + g^2 \phi^2 N^2/2$, has been considered in the literature mainly to discuss the production of heavy bosons needed for GUT baryogenesis [32]. Working in the Hartree approximation, it has been found [30] that a resonance is effective only provided the coupling g^2 satisfies $g^2 \gtrsim 10^{-7} (M_N/m_\phi)^4$. Taking into account all the other backreaction effects, a stronger lower bound on g has to be expected [25, 30], since the latter typically limits the growth of the fluctuations amplified by the resonance.

Very different bounds can be expected for different potentials. Consider for example $V(\phi) + (M_N + g\phi)^2 N^2/2$. In this case, due to the high initial amplitude of the inflaton oscillations, the total mass of N can vanish at some discrete points even for a coupling as small as $g^2 \sim 10^{-10} (M_N/m_\phi)^2$. Whenever $M_N + g\phi = 0$, parametric amplification of N occurs. Thus, the lower bound valid for the previous potential is considerably weakened. Although this second choice of the potential may seem *ad hoc*, we note that it is the one which arises in supersymmetric models if both the r.h. sneutrino mass and interaction with the inflaton are encoded in the superpotential, $W(N) \supset M_N N^2 + g\phi N^2$. We regard this as a very natural possibility.

The idea of a nonperturbative production associated to the vanishing of the total mass has been applied to leptogenesis in [33]. The analysis of [33] focused on the production of r.h. neutrinos, with a mass term of the form $(M_N + g\phi)\bar{N}N$. From the results of [33], and from the analytical computations of [34], it can be shown that a sufficient lepton asymmetry is generated if the mass of the r.h. neutrinos is higher

²A nonperturbative inflaton decay also occurs for hybrid inflation [27, 28, 29]. However, we will not discuss this possibility here.

will derive these bounds in section 2.3). Here we note that this high coupling can in principle destabilize through quantum effects the required flatness of the inflaton potential. This, in addition to the strong hierarchy between the r.h. neutrino mass and the electroweak scale, motivates the study of the supersymmetrized version of the mechanism proposed in [33].

One of our aims is to show explicitly that, in the supersymmetrized version of the above model, the nonperturbative production of the r.h. sneutrinos is much more efficient than the one of the neutrinos. Due to supersymmetry, the inflaton couples with the same strength both to the r.h. neutrinos and to the sneutrinos, so that if the former are produced at preheating this will also occur for the latter. However, while production of fermions is limited by Pauli blocking, the production of scalar particles at preheating is characterized by very large occupation numbers. This high production has typically a big impact on the dynamics of the inflaton field. The most immediate backreaction effect is the generation of an effective potential for the zero mode of the inflaton. This effective potential, taken into account in the Hartree approximation [25], is typically comparable with or even dominant over the tree level potential $V(\phi)$. There are however two equally important effects which are beyond the Hartree approximation. The first is due to the scatterings of the produced quanta against the zero mode of the inflaton. This destroys the coherence of the oscillations, thus ending the resonant production characterizing the early stage of preheating [25]. The second is the amplification of all the other fields to which the produced quanta are coupled. This is a very turbulent process, dominated by the nonlinear effects caused by the very high occupation numbers of the fields involved. As a result, all these mutually interacting fields are left with highly excited spectra far from thermal equilibrium [35]. Both these effects are denoted as *rescattering* [30].³

Rescattering strongly affects some of the outcomes of the analytical studies of preheating of bosons, which hardly go beyond the Hartree approximation. For this reason, the results presented in our work are obtained with numerical simulations on the lattice. More precisely, the code “LATTICEEASY” [36], by G. Felder and I. Tkachev, has been used (details are given in section 3). Full numerical calculations on the lattice are however rather extensive. We have found that the necessary computing time is reduced in the conformal case, that is with the inflaton potential $V(\phi) = \lambda \phi^4/4$, and with a r.h. sneutrino mass which is negligible during the early stages of preheating. For this reason, in our computations we fixed $M_N = 10^{11}$ GeV, which is smaller than the Hubble parameter during inflation, but still high enough to require a nonthermal production of the sneutrinos. The numerical results show a very efficient production of r.h. sneutrinos and inflaton quanta at preheating/rescattering. Even for a coupling inflaton-sneutrino as small as $g^2 \sim \text{few} \times 10^{-12}$, the produced quanta come to dominate the energy density of the Universe already within about the first 5 e-folds after the end of inflation. In particular, the energy density stored in sneutrinos is typically found to be a fraction of order one of the total energy density, so that a sufficient leptogenesis is easily achieved at their decay.

³In some works, the term *rescattering* refers only to the scatterings of the produced quanta against the zero mode of the inflaton. Here we keep the original meaning given in [30].

the superpotential term $h N H L \subset W$ (responsible for the Yukawa interaction which provides a Dirac mass to the neutrinos). Thus, one may expect that quanta of the latter fields are amplified by the rescattering of the r.h. sneutrinos produced at preheating. We study this possibility in section 4, showing that indeed the amplification occurs for a wide range of values of the coupling h . Part of the analysis follows the detailed discussion on rescattering given in [35], where the numerical code [36] used here was also employed. However, the analysis of [35] is focused on the production of massless particles, while we show that the non vanishing mass of the sneutrinos can have some interesting consequences. More precisely, when the sneutrino quanta become non relativistic (let us denote by $\hat{\eta}$ the time at which this happens) their rescattering effects become much less efficient. Thus, a strong amplification of the MSSM fields at rescattering can take place only if the coupling h is sufficiently large so that the amplification occurs before $\hat{\eta}$. As a consequence, for massive sneutrinos and for small values of h , the number of MSSM quanta produced at rescattering is an increasing function of h . However, the production is actually *disfavored* when the coupling h becomes too high. This is simply due to energy conservation, since the energy associated to the interaction term between the sneutrino and the MSSM fields cannot be higher than the energy initially present in the sneutrino distribution (equivalently, one can say that, for a too high coupling h , the non vanishing value of the sneutrinos gives a too high effective mass to the MSSM fields, which prevents them from being too strongly amplified). Posing quantitative bounds on the coupling h would require some better (analytical) understanding of the details of rescattering than we presently have. However, the numerical results shown in section 4 may give an idea of the expected orders of magnitude.

An important remark is in order. When we speak about the amplification of MSSM fields coupled to the r.h. sneutrinos we have actually in mind amplification of D -flat directions (let us generally denote them by X). Indeed, D -terms provide a potential term of the form $\Delta V \sim g_G^2 |Y|^4$ for any scalar non flat direction Y . Since g_G is a gauge coupling ($g_G = O(10^{-1})$), we expect such terms to prevent a strong amplification of Y , again from energy conservation arguments. Another important issue which emerges when gauge interactions are considered is whether gauge fields themselves are amplified at rescattering. We believe that, at least in the model we are considering, also the amplification of gauge fields will be rather suppressed. The scalar distributions amplified at rescattering break much of the gauge symmetry of the model. This gives the corresponding gauge fields an effective mass in their dispersion relation (analogous to the thermal mass acquired by fields in a thermal bath) of the order $m^2 \sim g_G^2 \langle X^2 \rangle$. As we extensively discuss in the paper, in the class of models we are considering the nonthermal distributions formed at rescattering are characterized by a typical momentum several orders of magnitude smaller than this mass scale. For this reason, one can expect that such heavy gauge fields cannot be strongly amplified.⁴ In our opinion, an explicit check of these conjectures by means of numerical simulations could be of great interest, especially considering the great importance that gauge fields could have for

⁴Gauge fields which are not coupled to the fields generated at rescattering will not acquire this high effective mass. However, being uncoupled, they will not be amplified either.

To conclude, we discuss the production of gravitino quanta from the scalar distributions generated at rescattering. We already mentioned that in order to avoid a thermal overproduction of gravitinos a lower bound has to be set on the reheating temperature T_{RH} of the thermal bath, $T_{RH} \lesssim \text{few} \times 10^{10}$ GeV [21]. The requirement of a low reheating temperature can be seen as the demand that the inflaton decays sufficiently late, so that particles in the thermal bath have sufficiently low number densities and energies when they form. If $H \simeq 10^{12}$ GeV at the end of inflation, and if the scale factor a is normalized to one at this time, the generation of the thermal bath cannot occur before $a \simeq 10^7$. Gravitino overproduction is avoided by the fact that in the earlier times most of the energy density of the Universe is still stored in the coherent inflaton oscillations. On the contrary, we have already remarked that preheating/rescattering lead to a quick depletion of the zero mode in the first few e-folds after the end of inflation.⁵

The question whether also the distributions formed at rescattering may lead to a gravitino problem is thus a very natural one, and section 5 of the paper is devoted to some considerations on this regard.⁶ To provide at least a partial answer to this question, we distinguish the period during which rescattering is actually effective from the successive longer thermalization era. The computation of the amount of gravitinos produced during the earlier stages of rescattering appears as a very difficult task. The numerical simulations valid in the case of bosonic fields indicate that a perturbative computation (with dominant $2 \rightarrow 2$ scatterings taken into account) can hardly reproduce the numerical results, and that probably $N \rightarrow 2$ processes ($N > 2$) have also to be taken into account (we discuss this point in more details in section 4). It is expected that the same problem will arise also for the computations of the quanta of gravitinos produced by the scalar distributions which are being forming at this stage. The end of rescattering/beginning of the thermalization period is instead characterized by a much slower evolution of the scalar distributions. In particular, the total occupation number of all the scalar fields is (approximatively) conserved, which is interpreted [35] by the fact that $2 \rightarrow 2$ processes are now determining the evolution of their distributions. Motivated by this observation, we assume that $2 \rightarrow 2$ interactions are also the main source of production for gravitinos from this stage on. In the thermal case, the gravitino production is dominated by processes having a gravitationally suppressed vertex (from which the gravitino is emitted) and a second vertex characterized by a gauge interaction with one outgoing gaugino. However, we believe that in the present context these interactions will be kinematically forbidden, due to the high effective mass-squared that gauginos acquire from their interaction with the scalar distributions (the argument follows the one already given for gauge fields). Once again we notice that the system is still effectively behaving as a condensate: the number densities of the scalar distributions are set by the quantity $\sqrt{\langle X^2 \rangle}$, which is much higher than the typical momenta of the distributions themselves. This generates a high effective mass for all the particles “strongly” coupled to these scalar fields. A further comparison with the case of a thermal distribution may be useful: in the latter case both the typical

⁵The situation is even more enhanced for hybrid inflationary model, in which the energy density of the zero modes of the scalars gets dissipated within their first oscillation [27, 28, 29].

⁶We acknowledge very useful discussions with Patrick B. Greene and Lev Kofman on this issue.

mechanism	N Yukawa h	N mass	ϕ - N coupling
thermal	$10^{-5}\text{eV} < \tilde{m}_1 < 10^{-3}\text{eV}$	$10^9\text{GeV} \lesssim M_1 \lesssim T_{RH}$	irrelevant
Affleck–Dine	$10^{-9}\text{eV} < m_{\nu_1} < 10^{-4}\text{eV}$	$M_i < H_{\text{infl}}$	$\begin{cases} M_i^{\text{eff}} < H_{\text{infl}} \\ (M_i^{\text{eff}})^2 < 0 \end{cases}$
$\left. \begin{array}{l} \text{pert.}\phi \\ \text{decay} \end{array} \right\}$	$\Gamma_{LV} < H(\tau_i)$	$\begin{cases} M_i < m_\phi/2 \\ M_i > m_\phi/2 \end{cases}$	$BR(\phi \rightarrow N_i N_i) \sim 1$
N preheating eq. (15)	$\Gamma_{LV} < H(\tau_i)$	$M_i \gtrsim 10^{14} \text{ GeV}$	$g_i \gtrsim 0.03$
\tilde{N} preh./resc. eq. (20)	$\Gamma_{LV} < H(\tau_i)$	$M_i \lesssim g_i 10^{17} \text{ GeV}$	$g_i \gtrsim \sqrt{\lambda}$

Table 1: Summary of parameters for which leptogenesis could work, for different r.h. (s)neutrino production mechanisms. In (s)neutrino production mechanisms which do not require the Yukawa coupling (non-thermal mechanisms), the constraint on the Yukawa matrix is that lepton number violating interactions in the thermal soup be out of equilibrium after the r.h. (s)neutrinos decay at τ_i . This also implies $M_i > T(\tau_i)$, and possibly additional constraints on L violating processes mediated by $N_j, j \neq i$. Recall \tilde{m}_i parametrises the N_i decay rate, and is defined after eqn. (13). The Affleck–Dine mechanism proceeds through generation of large expectation values either for a small [65] or a tachyonic [23] effective mass of sneutrinos M^{eff} during inflation. The asymmetry made by the perturbative decay of the inflaton can be generated by the on-shell r.h. (s)neutrinos [54], which subsequently decay, or by the decay via off-shell r.h. (s)neutrinos (N_i^*) to Higgses and leptons [55]. The properties of nonperturbative (s)neutrino production analysed in the present paper is summarized in the last two lines (λ is the inflaton self-coupling). Other scenarios for nonperturbative production after the end of inflation can be envisaged, with model-dependent results.

temperature of the system. As should be clear from the above discussion (see also [35]), the thermalization of the distributions produced at rescattering necessarily proceeds through particle fusion. Only after a sufficiently prolonged stage of thermalization, will the system be sufficiently close to thermodynamical equilibrium so to render processes as the one discussed above kinematically allowed.

In section 5 we show that if this class of processes is indeed kinematically suppressed, the production of gravitinos from the distributions formed at rescattering is sufficiently small. However, we remark that this analysis still leaves out the gravitino production which may have occurred at the earlier stages of the rescattering period. Whether this production may be sufficiently strong to overcome the limits from nucleosynthesis remains an open problem.

Let us finally summarize the plan of the paper. In section 2.1 we introduce our notation and briefly discuss some neutrino low energy phenomenology. In section 2.2 we discuss leptogenesis with a thermal production of the r.h. (s)neutrinos. Leptogenesis with a nonthermal production of r.h. neutrinos is reviewed in the following subsection. The supersymmetric version of this model is presented in section 3, where we study the nonthermal production of sneutrino quanta. Section 4 is devoted to the amplification of the MSSM D -flat directions due to the rescattering of the r.h. sneutrino quanta. The discussion on the gravitino production is presented in section 5, apart from a few technical details which can be found in the appendix.

2 See-Saw Phenomenology and Leptogenesis

In this section, we introduce our notation for the SUSY see-saw and outline its low-energy implications. The aim is to make contact between realistic see-saw models, and the one generation toy models in which we will study the sneutrino production. We discuss the lepton asymmetry that can be produced in (s)neutrino decay, which implies a lower bound on the mass of the lightest r.h. (s)neutrino. Then, we briefly review different mechanisms for r.h. neutrino production, namely thermal and non thermal. The terms neutrino and sneutrino will be used interchangeably in discussing thermal production, which is similar for bosons and fermions. Concerning the nonthermal case, instead, different results are obtained for the two species, and in section 2.3 we review the ones for the neutrinos. Nonthermal production of sneutrinos is instead discussed in the next section.

2.1 General Framework

Let us consider the Minimal Supersymmetric Standard Model (MSSM) extended with three r.h. neutrino superfields N_i (sometimes called the minimal supersymmetric see-saw model). The relevant couplings of the r.h. neutrinos are given by the superpoten-

$$W_N = h_{ji} L_i \cdot H_u N_j + \frac{1}{2} M_k N_k^2, \quad (1)$$

where L_i and H are the lepton and the Higgs doublets, respectively, and h is a 3×3 complex Yukawa matrix. We will neglect the phases in our analysis of N production, because CP violation is not required for this process. We work in the r.h. neutrino mass basis, where the mass matrix M is diagonal, and we disregard the possibility of nearly degenerate r.h. (s)neutrinos [37, 38, 11] (i.e. the difference of neutrino masses is of order their mass).

The lepton asymmetry produced in the decay of N_i can be written

$$Y_L \equiv \frac{N_L - N_{\bar{L}}}{s} = \epsilon_i \frac{N_{N_i}}{s} \kappa_i, \quad (2)$$

where N_{N_i} is the total number density of the i th heavy (s)neutrino species prior to its decay, s is the entropy density at decay ⁸, κ_i parametrises washout effects due to subsequent lepton number violating interactions, and ϵ_i arises from the CP violation of the N_i decay. It is given by [38]

$$\begin{aligned} \epsilon_i &\equiv \frac{\Gamma(N_i \rightarrow L H) - \Gamma(N_i \rightarrow \bar{L} \bar{H})}{\Gamma(N_i \rightarrow L H) + \Gamma(N_i \rightarrow \bar{L} \bar{H})} \\ &= \frac{1}{8\pi (h h^\dagger)_{ii}} \sum_j \text{Im} \left[(h h^\dagger)_{ij}^2 \right] f\left(\frac{M_j^2}{M_i^2}\right), \end{aligned} \quad (3)$$

where $f(x) = \sqrt{x}[2/(x-1) + \ln(1/x+1)]$ for hierarchical r.h. neutrino masses.

We suppose for the moment that some number density of N_i is produced in the early Universe, and concentrate on how large an asymmetry can be generated. The asymmetry ϵ_i is determined by the masses and couplings of the r.h. (s)neutrinos, which are given in eqn. (1). However, it can be related to, and therefore constrained by, low energy observables.

The CP asymmetry produced in the decay of a r.h. (s)neutrino can conveniently be parameterized as

$$\epsilon_i = \frac{3}{8\pi} \frac{M_i m_3}{\langle H \rangle^2} \delta_{CP} \simeq 10^{-6} \left(\frac{M_i}{10^{10} \text{ GeV}} \right) \left(\frac{m_3}{0.05 \text{ eV}} \right) \delta_{CP}. \quad (4)$$

By using eqs. (6) and (3), it is possible to show [13, 14, 15] that for the case ϵ_1 , δ_{CP} satisfies the upper bound ⁹

$$|\delta_{CP}| \leq 1. \quad (5)$$

By combining the two last expressions, one finds an upper bound on the parameter ϵ_1 which scales linearly with the r.h. (s)neutrino mass M_1 . We will shortly see that this implies a lower bound on M_1 for leptogenesis to be viable.

⁷The superfield N written in eqn. (1) actually denotes an anti-(s)neutrino. However, for brevity we will refer to it as a (s)neutrino.

⁸Any subsequent entropy production leads to further dilution of the asymmetry.

⁹We will use the parametrisation (4) for all the ϵ_i , $i = 1..3$. It is possible that $\delta_{CP} \leq 1$ for ϵ_2 and ϵ_3 (assuming no cancellations in the formulae), although this has not been shown.

The light neutrino mass matrix is obtained by integrating out the heavy r.h. neutrinos to give the see-saw formula

$$m_\nu = -h^T M^{-1} h \langle H_u^0 \rangle^2. \quad (6)$$

We will assume that the light neutrino masses m_i are hierarchical, so $m_3 \simeq \sqrt{\Delta m_{atm}^2}$ [8].

If h is written in the charged lepton mass eigenstate basis (neutrino flavour basis), then m_ν is diagonalised by the MNS matrix U [39], which can be written $U = V \cdot \text{diag}(e^{-i\phi/2}, e^{-i\phi'/2}, 1)$, with

$$V = \begin{pmatrix} c_{13}c_{12} & c_{13}s_{12} & s_{13}e^{-i\delta} \\ -c_{23}s_{12} - s_{23}s_{13}c_{12}e^{i\delta} & c_{23}c_{12} - s_{23}s_{13}s_{12}e^{i\delta} & s_{23}c_{13} \\ s_{23}s_{12} - c_{23}s_{13}c_{12}e^{i\delta} & -s_{23}c_{12} - c_{23}s_{13}s_{12}e^{i\delta} & c_{23}c_{13} \end{pmatrix}. \quad (7)$$

In this matrix, $c_{12} = \cos \theta_{12}$, and so on. Atmospheric [8, 40] and solar [41, 42] data imply that θ_{12} and θ_{23} are large, approaching $\pi/4$. θ_{13} is constrained to be $\lesssim 0.1$ by the CHOOZ experiment [43].

In a supersymmetric scenario, there is additional information about h and M available in the slepton mass matrix. The neutrino Yukawa $h^\dagger h$ appear in the renormalization group equations for the soft slepton masses, and thereby induce flavour violating slepton mass terms [44]: $[\tilde{m}_L^2]_{ij}$. In a simple-minded leading log approximation, these off-diagonal mass matrix elements are

$$[\tilde{m}_L^2]_{ij} \simeq \frac{(3m_0^2 + A_0^2)}{8\pi} [V_L]_{ki}^* [V_L]_{kj} h_k^2 \log \left(\frac{M_k}{M_{GUT}} \right) \quad (8)$$

where h_i are the eigenvalues of h , m_0^2 and A_0 are soft parameters at the GUT scale, and we introduce a new matrix V_L which diagonalises $h^\dagger h$ in the charged lepton mass eigenstate basis ($V_L h^\dagger h V_L^\dagger = \text{diagonal}$). The branching ratio for $\ell_j \rightarrow \ell_i \gamma$ can be roughly estimated as [44]:

$$\text{BR}(\ell_j \rightarrow \ell_i \gamma) \propto \frac{\alpha^3}{G_F^2} \frac{|[\tilde{m}_L^2]_{ij}|^2}{\tilde{m}_L^8} \tan^2 \beta \quad (9)$$

where \tilde{m}_L^2 is the slepton mass scale. The experimental bound $BR(\mu \rightarrow e \gamma) < 1.2 \times 10^{-11}$ [45] implies $[\tilde{m}_L^2]_{\mu e} \lesssim 10^{-3} \tilde{m}_L^2$, for $\tilde{m}_L \simeq 100$ GeV. This constrains the angles in V_L for given neutrino Yukawas h_i .

It can be shown that M and h have the same number of parameters as the weak scale neutrino and slepton mass matrices. Furthermore, in a SUSY scenario with universal soft masses at the GUT scale, M and h can be parametrised with \tilde{m}_L^2 and m_ν [46]. The r.h. neutrino masses and Yukawa couplings can be therefore reconstructed (*in principle*, but not in practise [46]) from the weak scale neutrino and sneutrino mass matrices, so that ϵ_i can be expressed in terms of weak scale variables.

An analytic approximation for ϵ_1 can be found in [47, 48]:

$$\epsilon_1 \simeq \frac{3h_1^2}{8\pi \sum_j |W_{1j}|^2 m_j^2} \text{Im} \left\{ \frac{\sum_k W_{1k}^2 m_k^3}{\sum_n W_{1n}^2 m_n} \right\}, \quad (10)$$

is the rotation from the basis where the ν_L masses are diagonal to the basis where $h^\dagger h$ is diagonal.

h_1 is in practise unmeasurable; however, if h has a hierarchy similar to the up Yukawa matrix h_u , then $h_1^2 \sim 10^{-8}$, and ϵ_1 will only be large enough if there is some enhancement from the imaginary part.

There are two simple limits for the matrix W , which are motivated by model building. The first is $V_L \simeq 1$, and corresponds to an almost diagonal slepton mass matrix (in the charged lepton mass eigenstate basis). This means that the large mixing observed in the MNS matrix U must come from the r.h. sector [49]. The second option is $W \simeq 1$, so $V_L \simeq U^\dagger$. This would arise if the large ν_L mixing is induced in the l.h. sector [50].

In the $V_L = 1$ case, eqn. (10) gives [47]:

$$\epsilon_1 \simeq -\frac{3h_1^2}{8\pi D} \text{Im} \left\{ \frac{m_1^3 c_{13}^2 c_{12}^2 e^{i\phi} + m_2^3 c_{13}^2 s_{12}^2 e^{i\phi'} + m_3^3 s_{13}^2 e^{2i\delta}}{m_1 c_{13}^2 c_{12}^2 e^{i\phi} + m_2 c_{13}^2 s_{12}^2 e^{i\phi'} + m_3 s_{13}^2 e^{2i\delta}} \right\} \quad (11)$$

$$\simeq -\frac{3h_1^2}{4\pi} \left\{ \left(\frac{m_3}{m_2} \right)^3 2s_{13}^2 \sin(2\delta - \phi') - \frac{m_1}{m_2} \sin(\phi - \phi') \right\}, \quad (12)$$

where $D = m_1^2 c_{13}^2 c_{12}^2 + m_2^2 c_{13}^2 s_{12}^2 + m_3^2 s_{13}^2$, and in the second equation, the solar and atmospheric angles have been taken to be $\pi/4$. If we estimate the phases to be $O(1)$, $h_1 \sim$ the up Yukawa, and $m_3^2/m_2^2 \sim \Delta m_{atm}^2/\Delta m_{sol}^2$, this gives $\epsilon \lesssim 10^{-7}(s_{13}/.1)^2$, where we have scaled the unmeasured angle θ_{13} by its upper bound.¹⁰ This is barely large enough for thermal leptogenesis. However, we remind that h_1 is unknown and it can well be $h_1 > 10^{-4}$.

The second case, where $W \simeq 1$, can arise if M and $h^\dagger h$ are almost simultaneously diagonalisable¹¹. For small angles in W , the approximation for ϵ can be extracted from (11), replacing the angles of the MNS matrix by the angles of W , and setting the cosines $\rightarrow 1$. When $W \sim 1$, then $V_L \simeq U^\dagger$, so it is the MNS angles that appear in equation (8), and $BR(\mu \rightarrow e\gamma) < 1.2 \times 10^{-11}$ [45] implies an upper bound on the CHOOZ angle $\theta_{13} < .02$ (for $\tilde{m}_L = 100$ GeV, $h_3 = 1$)[51].

To conclude, we briefly comment on the parameter κ_i . After the asymmetry is generated in the out-of-equilibrium decay of the r.h. neutrino, lepton number violating interactions which could wash out the asymmetry must be out of equilibrium. This is a fairly straightforward requirement when considering the decay of the lightest r.h. neutrino N_1 [12]; it is more complicated in the case of $N_{2,3}$ decaying at $T \gtrsim M_1$ [52]. The fraction of the asymmetry which survives these interactions is $\kappa_i \leq 1$.

2.2 Thermal N_i production

We now consider the case where the lightest r.h. (s)neutrino N_1 is thermally produced after T_{RH} . With hierarchical r.h. neutrino masses, one can typically assume the lepton

¹⁰This estimate is not significantly changed if the angles in V_L are small compared to θ_{13} . CHOOZ experiment: $\theta_{13} \lesssim .1$. If $\theta_{13} \ll .1$, but $\theta_{13} < [V_L]_{12}, [V_L]_{13} < .1$, then the formula for ϵ is similar to (12), with the replacement $\theta_{13} \rightarrow \theta_{L1j}$ and $\delta \rightarrow \varphi_{1j}$ (where $[V_L]_{12} = \cos \theta_{L13} \sin \theta_{L12} e^{i\varphi_{12}}$, $[V_L]_{13} = \sin \theta_{L13} e^{i\varphi_{13}}$).

¹¹The ‘‘almost’’ is important; if $W = 1$, there is no CP violation, so $\epsilon = 0$.

shall see, this is a self-consistent assumption, because M_2 and M_3 will turn out to be larger than T_{RH} .

To generate a lepton asymmetry, the decay of the N_1 should proceed out of equilibrium. More quantitatively, the ratio of the thermal average of the N_i decay rate and of the Hubble parameter at the temperature $T \simeq M_i$,

$$\frac{\Gamma_{N_i}}{2H}|_{T=M_i} \equiv \frac{\tilde{m}_i}{2 \times 10^{-3} \text{ eV}} , \quad (13)$$

should be less than unity to have an unambiguously out-of-equilibrium decay. The parameter \tilde{m}_i is defined as $\Gamma_i \langle H \rangle^2 / M_i^2$, where $\langle H \rangle$ is the Higgs vev. However, \tilde{m}_1 cannot be taken too small if N_1 is produced thermally [12]. Indeed the quantity \tilde{m}_1 controls the strength of the interactions of the N_1 with MSSM degrees of freedom, and an efficient thermal production via Yukawa interactions typically requires $\tilde{m}_1 \gtrsim 10^{-5}$ eV.

To account for both these effects, $\Gamma_{N_1} \sim H(T = M_1)$ should be taken, so the decay is only barely out of equilibrium, and the final lepton asymmetry has to be computed by integrating the full set of relevant Boltzmann equations [10, 11]. These computations show that a significant portion of the lepton asymmetry is erased by lepton-number violating processes, and that only a fraction $\kappa \lesssim .1$ or less typically survives. Starting from N_1 in thermal equilibrium at $T > M_1$, and collecting all the above informations, the final baryon asymmetry can be estimated to be

$$Y_B \simeq 10^{-10} \left(\frac{M_1}{10^{10} \text{ GeV}} \right) \left(\frac{m_3}{0.05 \text{ eV}} \right) \left(\frac{\kappa}{0.1} \right) \delta_{\text{CP}} . \quad (14)$$

This expression has to be compared with the baryon asymmetry required by Big Bang Nucleosynthesis $Y_B \equiv N_B/s \sim (1.7-8) \times 10^{-11}$ [53, 3]. As we have anticipated, we see that thermal leptogenesis can be a viable mechanism if the mass of the r.h. neutrino N_1 is sufficiently high. From eqn. (14) we find the lower bound $M_1 \gtrsim (10^9 - 10^{10})$ GeV, although it is fair to say that higher values are required if the bound (5) is not saturated. If N_1 is generated from the thermal bath, a reheating temperature greater than M_1 is required. In the most favorable case, the required value is only marginally compatible with the bound imposed by gravitino overproduction from the thermal bath [21]. To overcome the potential conflict between the gravitino bound and the requirement of a reheating temperature high enough for leptogenesis, the possibility of producing left-handed neutrinos *non thermally* has been envisaged [33]. The following considerations will be focused on this framework.

2.3 Nonthermal production of right handed neutrinos

Many alternatives to thermal (s)neutrino production have been considered in the literature. We concentrate here on mechanisms that involve a direct coupling of N to the inflaton, although in the next section we will comment on differences and similarities with the Affleck-Dine mechanism. The strength of this interaction, relative to the inflaton coupling to other degrees of freedom, is a free parameter; for appropriate values,

r.h. (s)neutrinos will also depend on the evolution of the inflaton between the end of inflation and reheating. If the inflaton decays perturbatively, right-handed neutrinos with masses less than half the inflaton mass could be produced in the decay [54]. For heavier r.h. neutrinos, one can also envisage the possibility that a sufficient leptogenesis is generated in processes in which they mediate a perturbative inflaton decay [55]. In both cases, the final lepton asymmetry will be proportional to the branching ratio of the inflaton into (either on- or off-shell) neutrinos. A branching ratio of order one is typically required.

Right-handed neutrinos with masses greater than that of the inflaton can be produced at preheating, if their interaction with the inflaton is strong enough. The production of heavy fermions (sneutrinos are discussed in the next section) in an expanding Universe was first discussed in ref. [33] (fermionic production in the conformal case was first studied in [56]), where a direct Yukawa coupling to the inflaton ϕ was considered, and the simplest chaotic inflationary scenario with a massive inflaton, $V(\phi) = m_\phi^2 \phi^2/2$, $m_\phi \simeq 10^{13}$ GeV, was assumed. The relevant part of the lagrangian is

$$\mathcal{L}_{N,\phi} = \bar{N} (M + g \phi) N, \quad (15)$$

where N is any one of the r.h. neutrinos. We assume that only one r.h. neutrino generation plays an important role in the generation of a lepton asymmetry, and therefore we drop the r.h. neutrino generation index for the remainder of this section. The generalization of the following analysis to three generations is straightforward, at least as long as the r.h. neutrino–inflaton coupling matrix g is diagonal in the r.h. neutrino mass basis (otherwise, the formalism of [57] should be used).

After the end of inflation, the inflaton condensate ϕ oscillates about the minimum of its potential with amplitude of a fraction of the Planck mass $M_P \simeq 1.22 \cdot 10^{19}$ GeV. The total effective mass of the fermion $M + g \phi(t)$ varies non adiabatically in time, and this leads to a (non perturbative) production of quanta of N . In particular, fermion production at preheating occurs whenever the total effective mass crosses zero. As a consequence, fermions with a mass up to

$$M_{\max} \simeq 5 \left(\frac{q}{10^{10}} \right)^{1/2} \times 10^{17} \text{ GeV}, \quad q \equiv \frac{g^2 \phi_0^2}{4 m_\phi^2} \simeq 3 g^2 10^{10} \quad (16)$$

can be produced [33], irrespectively on the value of the reheating temperature of the thermal bath which is formed at later times.

The abundance of neutrinos produced at preheating has been computed analytically [34], and it is most conveniently given in terms of the ratio

$$\frac{N_N}{\rho_\phi} \simeq \frac{1}{10^{10} \text{ GeV}} \frac{1.4 \times 10^{-14} q}{M_{10}^{1/2}} \left[\log \left(1.7 \times 10^3 \frac{q^{1/2}}{M_{10}} \right) \right]^{3/2}, \quad (17)$$

where we have defined $M_{10} = M / (10^{10} \text{ GeV})$. The above formula is valid as long as the backreaction of the produced neutrinos on inflaton dynamics is negligible, as it turns out to be the case as long as $q \lesssim 10^8$ [33, 34]. For larger values of q , the effectiveness

a lower bound on N_N . In what follows, we will conservatively assume $q < 10^8$.¹²

For a massive inflaton, the ratio (17) is constant until the inflaton condensate decays. If neutrinos decay before reheating has completed, the resulting baryon asymmetry reads

$$\begin{aligned} Y_B &= \frac{8}{23} Y_{B-L} = \frac{8}{23} \left(-\epsilon N_N \frac{3}{4} \frac{T_{RH}}{\rho_\phi} \right) = \\ &= 4 \cdot 10^{-14} \frac{q}{10^8} M_{10}^{1/2} \frac{T_{RH}}{10^9 \text{ GeV}} \frac{m_{\nu_3}}{0.05 \text{ eV}} \delta_{CP} \left[\text{Log} \left(1.66 \cdot 10^3 \frac{q^{1/2}}{M_{10}} \right) \right]^{3/2}. \end{aligned} \quad (18)$$

The ratio of N_N to the entropy is constant after reheating, unless it decreased due to some subsequent entropy production¹³. For a given value of the parameter q , the baryon asymmetry (18) is maximized when the mass M has its largest possible value (16). For $T_{RH} \simeq 10^9 \text{ GeV}$, $m_{\nu_3} = 0.05 \text{ eV}$ and $\delta_{CP} = 1$, imposing the condition (16) on eqn. (18) it is possible to see that the generation of the observed baryon asymmetry requires g larger than about 0.03. In particular, if we assume $q \simeq 10^8$, then the observed baryon asymmetry can be obtained for a mass of the right-handed neutrino of the order of $10^{14} - 10^{15} \text{ GeV}$.

Production of r.h. neutrinos at preheating can generate a large enough lepton asymmetry from neutrinos with $M_N > T_{RH}$, so constraints on T_{RH} do not translate into bounds on M_N . We notice however that the large inflaton-neutrino coupling required ($g \gtrsim 0.03$) can in principle modify through quantum effects the small mass or self coupling parameters characterizing the inflaton potential. This constitutes a further motivation for considering supersymmetric models, as we do in the remaining of this work. We will see that the production of the scalar partners of the r.h. neutrinos can significantly affect some of the above considerations.

3 Supersymmetric see-saw and nonthermal production of right handed sneutrinos

As we have seen in the previous section, preheating can have very important consequences for leptogenesis through the production of r.h. neutrinos [33]. In supersymmetric extensions of the see-saw model, the production of the supersymmetric partners of the neutrinos is even more important. Due to supersymmetry, the inflaton couples with the same strength both to the r.h. neutrinos and to the sneutrinos, so that if the former are produced at preheating this will also occur for the latter. However, while production of fermions is limited by Pauli blocking, the production of scalar particles

¹²One should also consider the perturbative decay of the inflaton quanta. Comparing eqn. (17) with the number of r.h. neutrinos produced perturbatively in one inflaton oscillation (i.e. the typical timescale for preheating), one can however see that the perturbative production is subdominant when kinematically allowed.

¹³For instance, this is the case if the r.h. neutrinos lifetime is long enough for them to come to dominate the energy density in the Universe, after reheating has completed.

the production of r.h. sneutrinos can be expected to be more significant than the one of neutrinos, as the numerical results presented below confirm.

Production of particles at preheating gives very model dependent results; nevertheless, some general features can be outlined, and the whole process can be roughly divided into three separate stages. The first of them is characterized by a very quick amplification of the fields directly coupled to the inflaton (and of the inflaton field itself, in the case of a sufficiently strong self-interaction) to exponentially large occupation numbers [25]. Very rapidly, the system reaches a stage in which the backreaction of the produced quanta, customarily denoted as rescattering [30], plays a dominant role. In the case of parametric resonance, the scatterings of the quanta against the zero mode of the inflaton destroy the coherence of the oscillations, thus ending the resonant production characterizing the early stage of preheating [25]. An equally important backreaction effect is the amplification of all the other fields to which the produced quanta are coupled. This is a very turbulent process, dominated by the nonlinear effects caused by the very high occupation numbers of the fields involved. As a result, all these mutually interacting fields are left with highly excited spectra far from thermal equilibrium. The latter is actually achieved on a much longer timescale, through an adiabatic (slow) evolution of the spectra, which characterizes the third and final stage of the reheating process.¹⁴

The first stage of preheating is well understood. Particle production is computed in a semi-classical approximation (for a rather general formalism in the case of several coupled fields see [57]), and analytical solutions have been obtained in a broad class of models [24, 25, 62, 31, 56, 34]. Analytical approximations break down when nonlinear processes become dominant. However, the high occupation numbers of the scalar fields involved allow a classical study of the system. Indeed, in the limit of high occupation numbers quantum uncertainties become negligible, and quantum probabilities show a classical (deterministic) evolution [63]. The latter can be better computed by means of lattice simulations in position space [30, 36, 29], where all the effects of backreaction and rescattering are (automatically) taken into account. A detailed discussion of rescattering and of the approach to thermal equilibrium has been given in [35], where the code “LATTICEASY” [36], by G. Felder and I. Tkachev, has been used. The numerical results presented in this paper are also obtained with this code.

For numerical convenience, we consider a chaotic inflationary scenario with a quartic potential for the inflaton. More specifically, we focus on the superpotential¹⁵

$$W(\Phi, N) = \frac{\sqrt{\lambda}}{3} \Phi^3 + \frac{1}{2} \left(\sqrt{2} g \Phi + M \right) N^2 . \quad (19)$$

¹⁴The thermalization of this system is a very interesting issue, which however we do not discuss in this work - see [35] for a more detailed study. Since in this case thermalization proceeds via particle fusion, an important role may be played by three or five point vertices, which shorten perturbative estimates of the thermalisation timescale [58] (other recent discussions on thermalization can be found in [59, 60, 61]).

¹⁵To embed the system in a supergravity context while preserving a flat potential for the inflaton field, one may impose [64] a definite parity for the Kähler potential $\mathcal{K} = \mathcal{K}(\Phi - \Phi^*)$. Doing so, the inflaton is identified with the real direction of the scalar component of Φ , and supergravity corrections can be neglected.

inflaton vev at the end of inflation	$\phi_0 = Mp/3$
inflaton self-coupling (eqn. (19))	$\lambda \simeq 9 \times 10^{-14}$
neutrino-inflaton coupling (eqn. (19))	$\tilde{g} = g^2/\lambda$
neutrino Yukawa coupling (eqn. (27))	$\tilde{h} = h^2/\lambda \quad (\leftrightarrow \tilde{m} = \frac{h^2 \langle H_u \rangle^2}{M})$
X number density (eqn. (22))	$N_{c,X} = [\text{comov. num. den.}]/[\sqrt{\lambda}\phi_0]^3$
X “mass” (eqn. (21))	$m_{\text{eff},X}^2 = \left(a^2 \langle \frac{\partial^2 V}{\partial \phi_i^2} \rangle - \frac{a''}{a} \right) / [\sqrt{\lambda}\phi_0]^2$

Table 2: translation table between quantities in plots and superpotential parameters (we give the eqn. where the parameter is defined). Recall that η is the conformal time coordinate, $\eta = 0$ at the end of inflation, and subsequently the scale factor is $a(\eta)/a(0) \simeq \eta/2 + 1$.

The second term of W reproduces the lagrangian (15) for the r.h. neutrinos. We denote the scalar components of the inflaton and of the r.h. neutrinos multiplets with ϕ and N , respectively. To simplify the numerical computations, the imaginary components of the scalar fields will be neglected. Therefore, after canonical normalization, $\phi \rightarrow \phi/\sqrt{2}$, $N \rightarrow N/\sqrt{2}$, we consider the scalar potential ¹⁶

$$V_{\text{scalar}} = \frac{\lambda}{4} \phi^4 + \frac{1}{2} (g\phi + M)^2 N^2. \quad (20)$$

The size of the temperature fluctuations of the Cosmic Microwave Background sets $\lambda \simeq 9 \times 10^{-14}$, while the neutrino mass M as well as the coupling g to the inflaton are model dependent parameters. The case $M = 0$ is analyzed in details in [35].

In figure 1 we show the time evolution of the comoving number density and of the comoving effective mass of the two scalars ϕ and N , for the particular choice of the parameters $M = 10^{11}\text{GeV}$, $\tilde{g} \equiv g^2/\lambda = 200$. The effective mass is defined as

$$m_{\text{eff},\phi_i}^2 \equiv a^2 \langle \frac{\partial^2 V}{\partial \phi_i^2} \rangle - \frac{a''}{a}, \quad (21)$$

where a is the scale factor of the Universe, normalized to one at the end of inflation, prime indicates derivative with respect to the conformal time η , while $\langle \dots \rangle$ denotes average over the sites of the lattice. The term a''/a appears in eqn. (21) because we are considering minimally (rather than conformally) coupled scalars, and it vanishes in a radiation-dominated background. The comoving number density is defined as the

¹⁶We neglect the term quartic in N , subdominant with respect to the mass term $M^2 N^2$ during most of the preheating/rescattering period, as well as the mixed term $\propto \sqrt{\lambda} g \phi^2 N^2$, negligible with respect to the mixed term present in eqn. (20) for $g^2 \gg \lambda$.

One may also be worried that, if the right-handed sneutrino is charged under some gauge group (as it generally happens in grand unified models) with a gauge coupling g_G not much smaller than one, the corresponding D -term $\propto g_G^2 |N|^4 \subset V$ could prevent the amplification of N at preheating-rescattering. However, at least as long as $\langle N \rangle$ is smaller than the scale at which the gauge symmetry is broken, this term gets actually compensated by a shift of the (much heavier) field responsible for the breaking of the symmetry, and it is thus absent from the effective potential for N [65, 22].

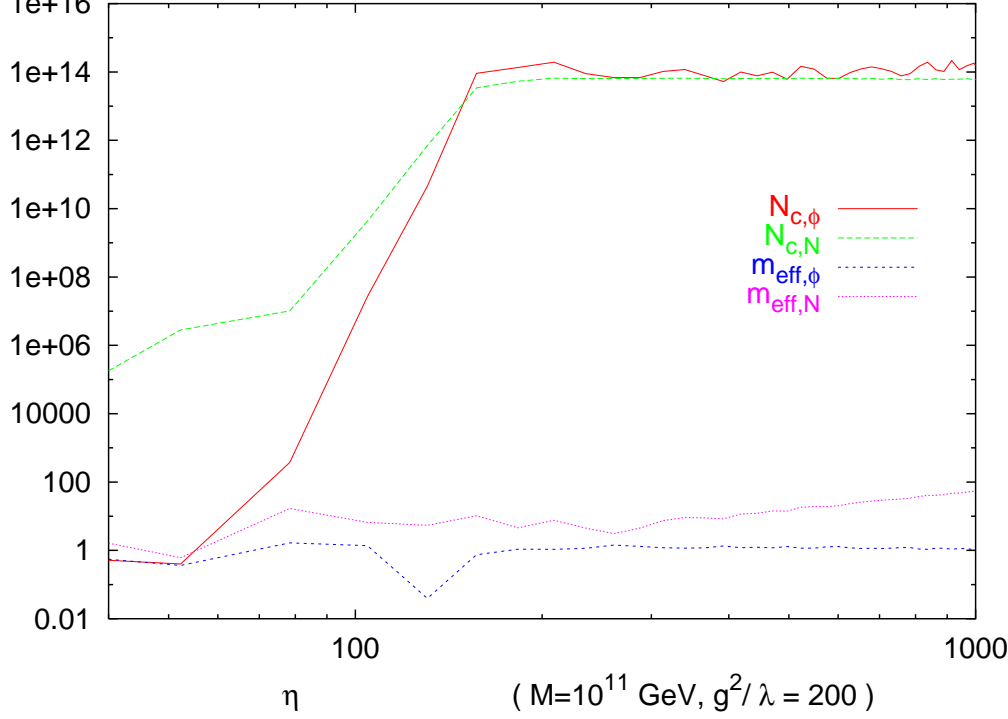


Figure 1: Time evolution of the comoving number density and of the effective masses of the inflaton ϕ and of the r.h. sneutrino N .

integral over momentum of the “occupation number”

$$\begin{aligned}
n_k(\eta) &\equiv \frac{1}{2} \left(\omega_k |f_k|^2 + \frac{1}{\omega_k} |f'_k|^2 \right), \\
\omega_k^2 &\equiv k^2 + m_{\text{eff}}^2,
\end{aligned} \tag{22}$$

where f_k denotes the Fourier transform (to be evaluated on the lattice) of the rescaled field $a\phi$. By definition [36], the quanta stored in the oscillating inflaton condensate do not contribute to $N_{c,\phi}$ in figure 1. The three quantities m_{eff} , N_c , and η are all shown in units of $\sqrt{\lambda}\phi_0 \simeq 1.25 \cdot 10^{12} \text{ GeV}$, with $\phi_0 \simeq 0.342 M_P$ denoting the value of ϕ at the end of inflation, to the appropriate power. All the numerical results presented in this work are obtained with a two dimensional lattice of size $L = 20 \left(\sqrt{\lambda}\phi_0 \right)^{-1}$ and with $N = 1024^2$ sites (see [36] for details).

Figure 1 exhibits the features that we have outlined at the beginning of this section, namely a quick stage of exponential growth of the occupation numbers followed by a period in which the occupation numbers are nearly constant. During the first stages of the process, the results presented reproduce very well the ones obtained in [35] for $M = 0$, since the “bare” mass of the r.h. neutrinos is initially negligible.

However, the presence of a non vanishing bare mass affects the subsequent evolution of the system. Indeed, when the value of the fluctuations in the sneutrino field becomes comparable with the amplitude of inflaton oscillations, the second term in eqn. (20) shifts the minimum of the effective potential of the inflaton, giving it an effective mass

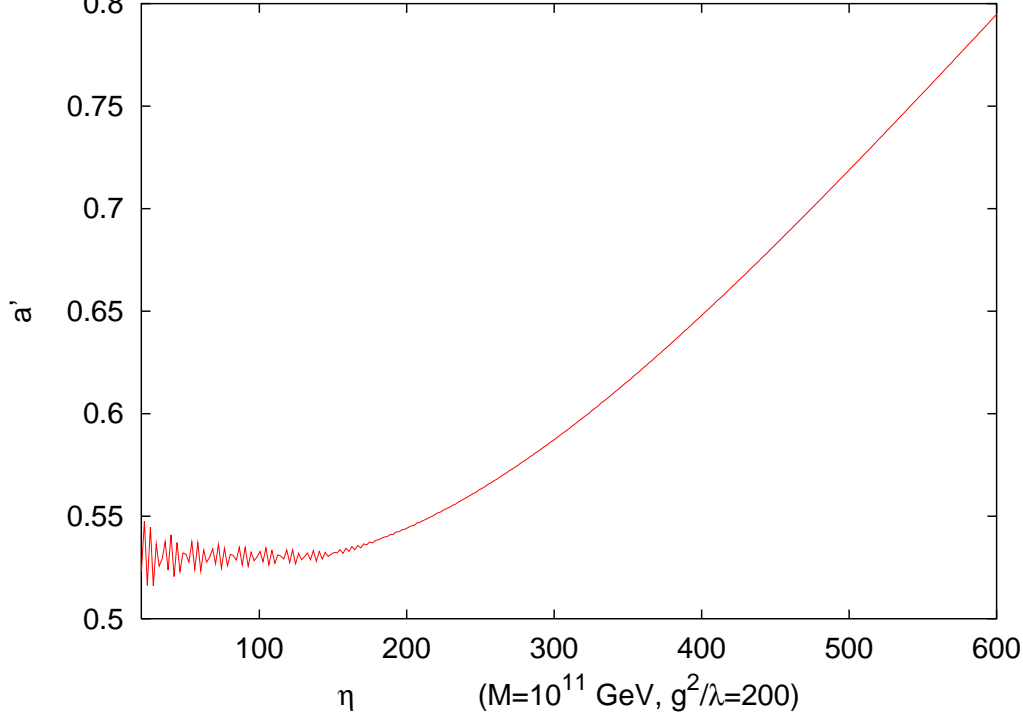


Figure 2: Time evolution of the derivative of the scale factor with respect to conformal time. A constant value indicates radiation domination, while a linear growth indicates matter domination.

that is roughly constant in comoving units. A stronger effect is related to the fact that the sneutrino itself is massive. In rescaled units and for the present choice of the parameters, we have

$$m_{\text{eff},N}^2 \simeq \left\langle \left(0.08 a + 14.1 \frac{a \phi}{\phi_0} \right)^2 \right\rangle. \quad (23)$$

Numerical results show that $\langle \phi \rangle \sim \phi_0/a$ (as one could also see by inspecting the potential (20) in Hartree approximation), so that the part in the above equation that depends on ϕ remains of order one. It follows that the r.h. neutrino mass M dominates the effective mass (23) for $\eta \geq \hat{\eta} \sim 350$, as clearly indicated by the growth of $m_{\text{eff},N}$ visible in figure 1 for $\eta > \hat{\eta}$.

Production of sneutrinos at preheating in the present model is strictly related to the production of fermions we have analysed in section 2.3. In particular, production occurs whenever the effective neutrino mass (23) crosses zero. Numerical results show that preheating is terminated by rescattering effects when the scale factor a is of the order of $a_{\text{resc}} \simeq 100$. As a consequence, sneutrinos with a bare mass up to $g \phi_0/a_{\text{resc}} \simeq g \cdot 10^{17} \text{ GeV}$ will be efficiently produced at preheating and will constitute a sizable fraction of the background energy¹⁷.

Numerical results show that after the onset of rescattering, the energy density

¹⁷Non-adiabatic production of sneutrinos can occur for bare masses as large as $g \phi_0/4 \sim g \cdot 10^{18} \text{ GeV}$, but the efficiency of the process will be much lower, because redshift effects will terminate the resonance before rescattering sets in.

in general large couplings correspond to a large interaction energy, and therefore to a smaller number density during the rescattering/thermalization stage [30, 35]. In particular, for this reason a large quartic self-coupling $g_N^2 |N|^4$ for the sneutrino would prevent it from getting large occupation numbers, since energy conservation would impose $\langle N^2 \rangle \propto \sqrt{\lambda/g_N^2}$. Numerical results also indicate that soon after the beginning of rescattering most of the r.h. neutrino quanta have a momentum of the order $k_* \sim 15 \sqrt{\lambda} \phi_0$.¹⁸ Thus, most of the r.h. neutrinos become non relativistic at a time not much greater than $\hat{\eta}$. From this time on, the energy density of the system redshifts as the energy density of matter. The transition between the two stages of matter and radiation domination is clearly visible in figure 2, where we show the time evolution of the derivative of the scale factor with respect to conformal time. As long as the neutrino mass is negligible, the energy density of the system redshifts as the one of radiation [31, 35], and the evolution of the scale factor is very well approximated by $a \simeq (1+t)^{1/2} \simeq \eta/2+1$, where we have set $t = \eta = 0$ at the end of inflation. Therefore, during the initial stage of radiation domination, a' is constant. In the following matter dominated stage $a \propto \eta^2$, and a' grows linearly with time.

To estimate the baryon asymmetry produced from the decay of the r.h. sneutrinos, we need to know the fraction of the entropy of the Universe that is generated in the decay. The baryon asymmetry will be [13]

$$\begin{aligned} Y_B &\simeq -\epsilon \frac{8}{23} \frac{N_N}{s_N} \frac{s_N}{s_{tot}} \simeq \frac{8}{23} \left(-\epsilon \frac{3}{4} \frac{T_N}{M} \right) \frac{s_N}{s_{tot}} \\ &\simeq 0.3 \times 10^{-10} \left(\frac{T_N}{10^6 \text{ GeV}} \right) \left(\frac{m_3}{0.05 \text{ eV}} \right) \frac{s_N}{s_{tot}} \delta_{\text{CP}} , \end{aligned} \quad (24)$$

where $s_N \propto g_*(T_N) T_N^3$ is the entropy produced in N decay, and s_{tot} is the total entropy of the Universe. If the r.h. sneutrinos dominate the Universe when they decay, then T_N is the temperature to which the Universe is reheated by the decay of the sneutrinos N , and $s_{tot} \sim s_N$. This condition is satisfied if the inflaton mainly decays only into one r.h. sneutrino family (as it is clearly the case in the one generation model we have studied numerically). However, if the inflaton couples to other scalars (also in particular to the other generations of sneutrinos), these could produce additional entropy.

At variance with the case of leptogenesis induced by the decay of right-handed neutrinos, analysed in section 2.3, the baryon asymmetry (24) does not depend on the r.h. (s)neutrino mass, that must only satisfy $M > T_N$ in order to prevent thermal regeneration of the r.h. sneutrinos after their decay. This is due the fact that, thanks to Bose statistics, r.h. sneutrinos can get large occupation numbers at preheating (whereas Pauli blocking makes fermion production less efficient), and they can easily represent a substantial fraction of the energy in the Universe.¹⁹

¹⁸The precise value of the typical momentum k_* of the distributions, as well as the time needed for N_c to saturate, are a nontrivial function of \tilde{g} , since different values of this parameter lead to different positions (in momentum space) and strengths of the resonance bands [31]. However, the rescattering stage destroys these resonance bands, making the dependence of k_* on \tilde{g} milder.

¹⁹Notice that for a massless inflaton, the inflaton energy redshifts as radiation, and non-relativistic

inflaton is obtained by considering the following superpotential, in the mass eigenstate basis for the r.h. neutrinos,

$$W = \frac{\sqrt{\lambda}}{3}\Phi^3 + h_{ji} L_i \cdot H_u N_j + \frac{1}{2}(M_j \delta_{jk} + \sqrt{2}g_{jk}\Phi)N_k N_j. \quad (25)$$

This gives a potential for the real components of the scalars

$$V = \frac{\lambda}{4}\phi^4 + \frac{1}{2}\sum_{i,j,k}(g_{ij}\phi + M_i\delta_{ij})(g_{ik}^*\phi + M_i\delta_{ik})N_j N_k + \dots \quad (26)$$

where dots include the terms involving the Yukawa h , which are not relevant for non-thermal N_i generation. We suppose for simplicity that the neutrino–inflaton coupling g_{ij} is diagonal. Energy considerations after rescattering [30, 35] lead to the expectation that the sneutrino family that is most strongly coupled to the inflaton is also the one that will have the smallest number density (clearly, on the assumption that all the sneutrinos are sufficiently coupled to be amplified). This is opposite to the scenario in which sneutrinos are produced by the perturbative decay of ϕ , where N_{N_i} is proportional to the inflaton branching ratio to N_i . However, the presence of the r.h. sneutrino bare masses, as well as the existence of couplings to the Standard Model degrees of freedom (whose effects will be analysed in detail in the next sections), can strongly affect these conclusions.

Although the resulting baryon asymmetry (24) has the same expression as the one reported in [13], the mechanism that led to a sneutrino dominated Universe is different from the generation of large expectation values considered in ref. [13] or in the Affleck–Dine [65] mechanism. Indeed, the latter is effective if during inflation the sneutrino (or, more generally, the amplified flat direction) has a mass much smaller than the Hubble rate. This requires (besides a sufficiently small bare mass $M \ll H$) that the amplified field does not get a large effective mass through its coupling to the inflaton. It is important to remark that the mechanism we are discussing can provide a sufficient leptogenesis even if the coupling g between the inflaton and the r.h. neutrino multiplets is much smaller than the one needed in non supersymmetric models, i.e. with only the production of the neutrinos taken into account, see eqn. (18). Anyhow, even couplings as small as $g^2 \sim 10\lambda$ prevent the formation of a large condensate during inflation. Therefore, the two mechanisms can lead to large occupation numbers in complementary regions of the parameter space.

Notice that the above discussion applies to every effective mass term that can arise in the potential for a (quasi) flat direction. In particular, it could be interesting to consider the effective mass of the order of the Hubble parameter that is generally induced by supergravity corrections [23], although in this case amplification effects may be weakened by the quick redshift characterizing the nonrenormalizable interactions.

neutrinos will easily dominate the energy in the Universe. If the inflaton is instead massive, then the energy in sneutrinos would in any case be a fraction of order unity of the background energy, and would start increasing after inflaton decay. As a consequence, the resulting baryon asymmetry would still be given (at most up to factors of order one) by the expression (24).

value will be generated during inflation [23], and the dynamics of preheating will turn out to be rather different from the one considered in the present section. This however requires a suitable nonminimal Kähler potential, and we will not discuss this possibility in this work.

Finally, it is worth stressing that both the leptogenesis scenario described in this section and the one considered in [13], although they are related to Affleck–Dine leptogenesis, are somehow different from it for what concerns the fulfillment of the Sakharov CP -violation condition [2]. In the Affleck–Dine scenario, the latter is achieved by the motion of the Affleck–Dine condensate (that requires coherence over many Hubble lengths), while in the mechanism we are analysing, CP -violating sneutrino decays are crucial in the generation of an asymmetry.

4 Production of light degrees of freedom at rescattering

The description presented above is further modified by the effects of the coupling of the r.h. neutrino multiplets to the l.h. leptons and Higgses, coming from the superpotential (1). Considering for simplicity only one generation, the superpotential (19) will be supplemented by

$$\Delta W = h N L \cdot H . \quad (27)$$

The resulting scalar potential contains several interaction vertices coming from F-terms. In addition, there are quartic contributions from the D-terms. The presence of the latter in the potential plays an important role for our discussion, since rescattering mostly affects the D-flat directions. To see this, consider the case in which the left-handed selectron and the charged scalar Higgs vanish. The D -terms for the neutral Higgs and sneutrino then take the familiar form of the tree level MSSM Higgs potential, with H_d^0 replaced by $\tilde{\nu}_L$

$$V_D = \frac{g_{\text{SU}(2)}^2 + g_Y^2}{8} \left| |H^0|^2 - |\tilde{\nu}_L|^2 \right|^2 . \quad (28)$$

The directions that are not D-flat (i.e. the ones for which $|\tilde{\nu}_L| \neq |H^0|$ in the present example) are characterized by a large (gauge) quartic coupling in the scalar potential. Due to this coupling, they cannot be significantly amplified by rescattering effects. On the contrary, D-flat directions have only quartic couplings coming from F-terms as $|\partial W / \partial N|^2$, whose strength h^2 will be typically taken $\ll 1$ in all the cases considered below. Indeed these quartic interactions among the D-flat directions will be neglected in our computation, since they can be relevant only during the thermalization stage, when the variances of these fields have grown to be sufficiently large.

The most important F-terms arising from the total superpotential are of the form $\sim h^2 |N|^2 (|\tilde{\nu}_L|^2 + |H^0|^2)$. If we denote by X the D-flat combination $|\tilde{\nu}_L| = |H^0|$, then the relevant coupling of X to the sneutrino field will be simply given by $h^2 N^2 X^2$ (as in the previous section, we consider for numerical convenience only real directions). Besides

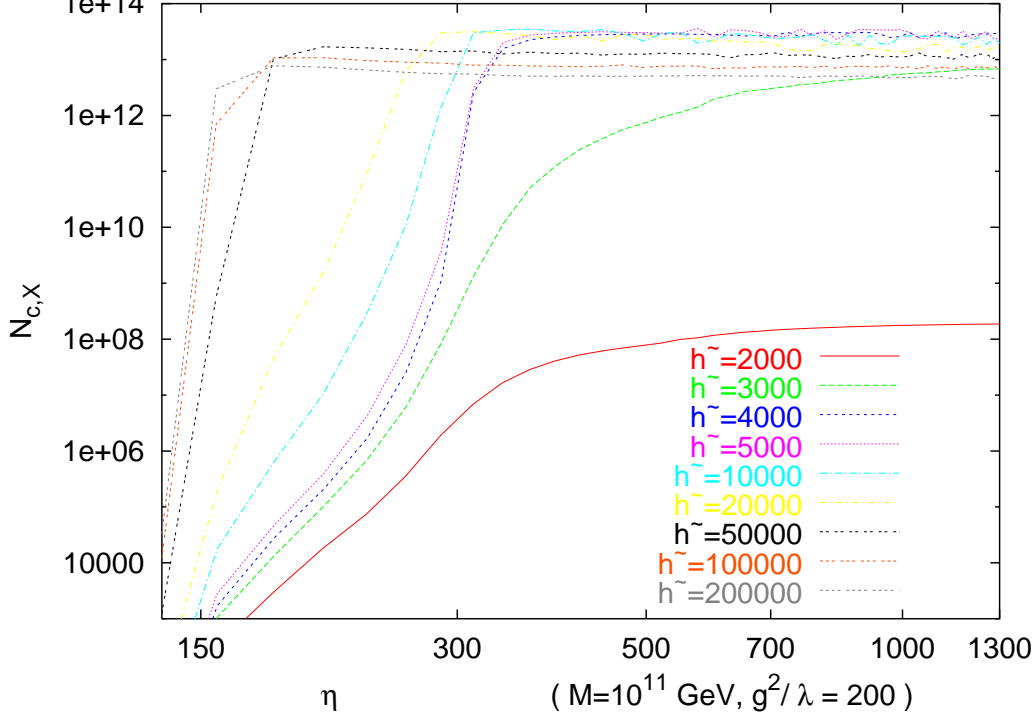


Figure 3: Time evolution of the comoving number density of the light quanta X .

the quartic term $\propto h^2 X^4$, we will also neglect the interaction term $\propto (g\phi + M) N X^2 \simeq M N X^2$, which is responsible for the late time decay of the r.h. sneutrinos (that is, the supersymmetric counterpart of the vertex which gives the decay of the r.h. neutrinos into Higgses and leptons).

We thus consider the simplified model characterized by the scalar potential

$$V(\phi, N, X) = \frac{\lambda}{4} \phi^4 + \frac{1}{2} (M + g\phi)^2 N^2 + \frac{1}{2} h^2 N^2 X^2. \quad (29)$$

Neglecting the imaginary directions of the scalar fields, as well as many of the interaction terms, allows a considerable reduction of the computing time needed for the numerical simulations (this is particularly welcome for the extensive computation that we discuss in the next section). The above discussion leads us however to believe that the simplified model should well describe the main features of the preheating and rescattering process for the supersymmetrized see-saw model with a nonperturbative production of the r.h. neutrinos.

In figure 3 we show the time evolution of the comoving number density of the quanta of X . As in the previous section, we have fixed $M = 10^{11}$ GeV, $\tilde{g} = 200$, while different values of the parameter $\tilde{h} \equiv h^2/\lambda$ are shown. Even if in the simplified model (29) the X field is not directly coupled to the inflaton, we see that (for suitable values of the coupling h) it can be highly amplified by the rescattering of r.h. sneutrino quanta.

Figure 3 shows that the growth of number of X particles is roughly exponential in time. When the effective sneutrino mass is varying non-adiabatically in time and is not

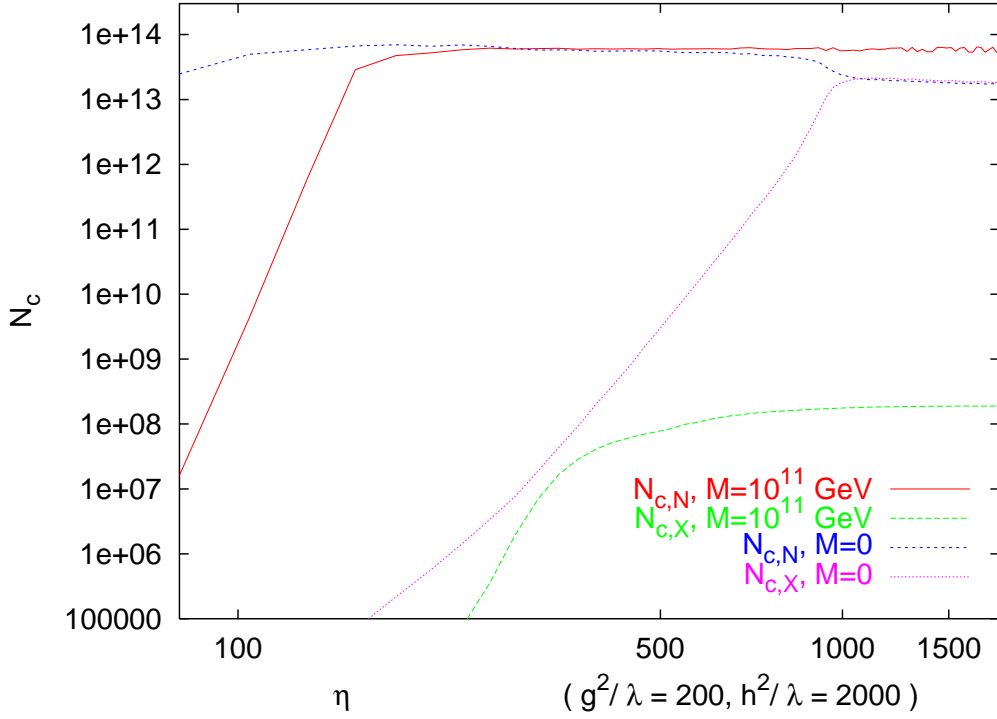


Figure 4: Comparison of the growth of the occupation numbers of N and X with and without the r.h. neutrino mass term. Notice that the amplification of X considerably weakens when the r.h. neutrino quanta become non-relativistic, $\eta \sim 350$.

negligible with respect to sneutrino typical momenta, the production of X particles cannot be analysed in terms of scatterings of sneutrinos. However, after the end of the parametric resonance period and the onset of rescattering, one can expect that a particle-like picture can give some description of the behavior of the system [35]. In this case, if the dominant contribution to X production process were given by the $2 \rightarrow 2$ scattering $NN \rightarrow XX$, the rate of growth of N_X should be proportional to h^4 . The lattice results appear to present a milder dependence on h , suggesting that the $NN \rightarrow XX$ scatterings alone cannot account for the production of X states. In a naive perturbative analysis, the contribution, from $(m+2) \times N \rightarrow XX$ processes, to the rate of growth of N_X , is proportional to the $m \times N \rightarrow XX$ rate times a factor roughly given by $h^4 N_N^2 / (4\pi p^6)$, where $p \sim 15\sqrt{\lambda}\phi_0/a$ is the typical momentum exchanged. Due to the high density of sneutrinos, the expansion parameter $h^4 N_N^2 / (4\pi p^6)$ is of order unity for the values of h we are considering. Therefore, strongly turbulent processes with many incoming sneutrinos can contribute significantly to the rate of growth of N_X . This is confirmed by the fact that the total number of particles decreases during the stage of generation of the X states, thus showing that particle fusion processes are dominant at this stage [35].

The main features shown in figure 3 are shared by the other evolutions with different \tilde{g} that we will consider in the next section, and they can be understood at least qualitatively. As could have been easily guessed, the timescale for the growth of $N_{c,X}$ is a decreasing function of \tilde{h} . We also notice from the figure that the amplification

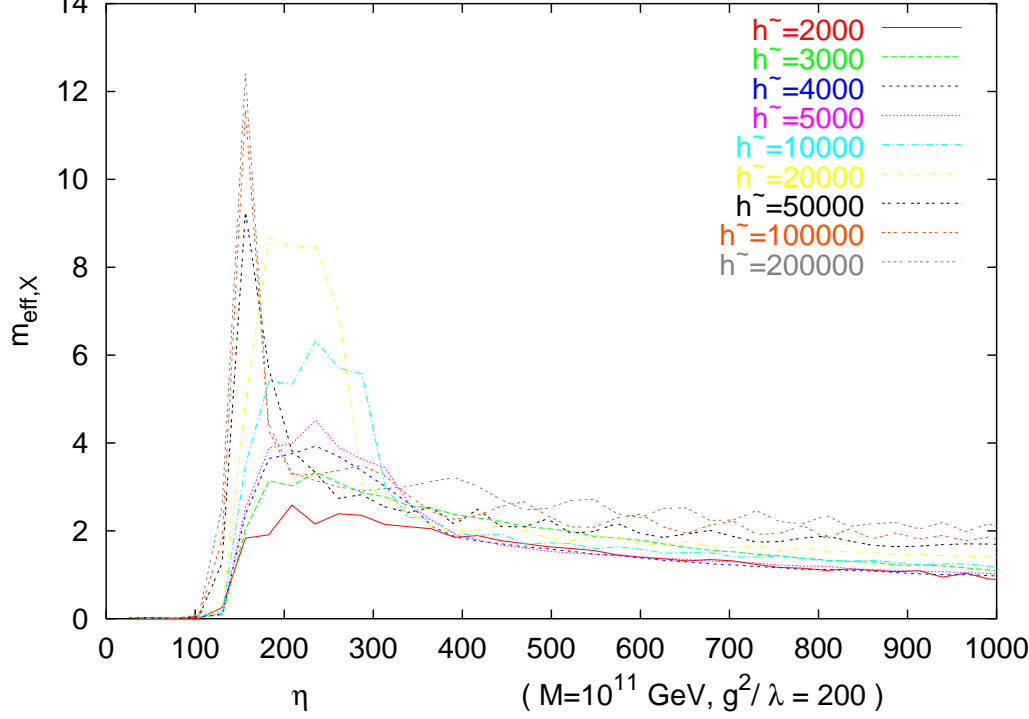


Figure 5: Time evolution of the effective mass of the quanta of X .

of X becomes less efficient as the quanta of N become non relativistic, at $\eta \sim 300$. This can be seen explicitly in figure 4, where we show the effect of the r.h. neutrino mass term on the growth of the comoving occupation numbers of the N and X fields. If the two fields are both massless, the rescattering of the quanta of N lifts the X to (practically) the same amplification [35]. We observe that the situation is completely different for the case in which the quanta of N have a sufficiently high mass. Indeed, when the amplification of X from the r.h. neutrino quanta substantially decreases when the latter become non-relativistic. As a consequence of the two effects mentioned in this paragraph, the asymptotic value of $N_{c,X}$ (at least in the time interval we have considered) decreases at small \tilde{h} .

Figure 3 also shows a decrease of the asymptotic $N_{c,X}$ for high \tilde{h} . As discussed in [35], for sufficiently high coupling h the two fields have comparable occupation numbers, $N_{c,N} \simeq N_{c,X}$. Clearly, the higher the coupling is, the sooner this (approximate) equipartition is reached. For a very high h , the potential energy associated to the last term of (29) then disfavors the production of the quanta of the two fields, in the same way as a high quartic coupling $\propto N^4$ added to the potential (20) would have prevented the amplification of the r.h. neutrinos in the two fields case.

In figure 5, we show the time evolution of the effective mass of the quanta of X . Comparing it with their distribution in momentum space, one realizes that most of the quanta are always in a relativistic regime.

The light quanta X generated at rescattering can in turn be responsible for the production of unwanted relics such as gravitinos. If unstable, gravitinos with a mass of the order of the electroweak scale (which is what we expect in models of gravitationally mediated supersymmetry breaking) disrupt the successful predictions of primordial nucleosynthesis, unless their abundance is below the very stringent bound [19]

$$Y_{3/2} \lesssim 10^{-14} (\text{TeV}/m_{3/2}) \quad . \quad (30)$$

In inflationary theories, several sources of gravitino production have been considered. The most standard of them is the perturbative production from the thermal bath formed at reheating. In this case, the above limit (30) translates into the upper bound $T_{\text{rh}} \lesssim \text{few} \times 10^{10} \text{ GeV}$ on the reheating temperature [21]. Other sources of gravitino production can be studied. Since gravitinos arise in supersymmetric models, a very natural “candidate” channel for their production is the decay of the inflaton into its supersymmetric partner (the inflatino) plus a gravitino. This process is however either kinematically forbidden [18] or strongly suppressed [66] by the fact that the difference between the inflaton and the inflatino mass is governed by the supersymmetry breaking scale, which is of the same order of the gravitino mass. The resulting gravitino production is sufficiently small. Recently, the generation of gravitinos at preheating has been extensively discussed, both concerning the relatively easier case of the transverse component [67], and the more delicate issue of the longitudinal component [68, 57, 69]. These studies are focused on the nonperturbative amplification of the gravitino field due to the coherent oscillations of the inflaton, and this mechanism of production is found to be sufficiently limited [57] provided that the inflationary sector of the theory is weakly (e.g. gravitationally) coupled to the one responsible for the present supersymmetry breakdown.

In this section we discuss a different possible source of gravitinos, namely a perturbative production from the nonthermal distributions of light MSSM quanta generated at rescattering. A comparison with the standard thermal production may be used as an initial motivation. Concerning the latter, the requirement of a low reheating temperature can be seen as the demand that the inflaton decays sufficiently late, so that particle in the thermal bath have sufficiently low number densities and energies when they form. If $H \simeq 10^{12} \text{ GeV}$ at the end of inflation, and if the scale factor a is normalized to one at this time, the generation of the thermal bath cannot occur before $a \simeq 10^7$. Gravitino overproduction is avoided by the fact that in the earlier times most of the energy density of the Universe is still stored in the coherent inflaton oscillations. As we have seen in the previous sections, this last assumption is not any longer valid if preheating and rescattering effects are important. Indeed, in the model considered above the energy density of the scalar distributions becomes dominant already when the scale factor is of order 100 (the precise value being a function of the parameters of the model).²⁰

²⁰The situation is even more enhanced for hybrid inflationary model, in which the energy density of the zero modes of the scalars gets dissipated within their first oscillation [27, 28, 29].

of gravitino production in the context of rescattering is certainly more difficult than the usual thermal production. While in the latter case a perturbative approach can be adopted, and the final result can be readily estimated by computing the rate of $2 \rightarrow 2$ processes with one gravitino in the final state, rescattering is a highly nonlinear phenomenon. In the bosonic case, we already remarked that naive perturbative estimates poorly reproduce the initial amplification of the fields X . Only towards the end of the rescattering stage the number densities of the amplified fields become sufficiently small so that $2 \rightarrow 2$ processes become dominant, as the (approximate) conservation of the total comoving occupation number at relatively late times signals [35]. Unluckily, Pauli blocking forbids fermionic fields to behave classically (in the sense discussed in section 3), and lattice simulations cannot be used. However, we may take the numerical results for bosons as a guideline. Also for the production of fermions more complicated processes than just $2 \rightarrow 2$ interactions could be relevant during most of the rescattering stage, while they should be subdominant at sufficiently late time. The latter is presumably set by the same timescale at which rescattering is seen to end in the numerical simulations described in the previous sections. With this in mind, we proceed to an estimate of the amount of gravitinos produced by $2 \rightarrow 2$ processes with the fields amplified at preheating/rescattering in the incoming state. We stress once more that this estimate can be reliable only from the beginning of the thermalization stage on, so it should provide a lower bound on the total number of produced gravitinos. It is possible that a higher amount of gravitinos is produced at earlier times, when nonlinear effects cannot be neglected.

This computation has been carried out in appendix, where also some details (i.e. concerning the quantization of the system and the choice to focus on the transverse gravitino component) are reported. As for the thermal production, the dominant $2 \rightarrow 2$ processes have only one gravitino in the final state, and hence only one gravitationally suppressed vertex. In the thermal case, the dominant processes have a gauge interaction as the second vertex, consider e.g. the process $X X \rightarrow \tilde{z} \psi_{3/2}$ with one higgsino in the propagator. For the nonthermal distributions of scalars that we are considering, however, such processes are kinematically forbidden. This is a crucial point, which poses a significant limit on the estimated production of gravitinos. We can easily understand it using the specific process just mentioned as an example: either the Higgses are not amplified at rescattering (so the above process is irrelevant) or the non-vanishing $\langle H^2 \rangle$ provides an effective mass to the zino produced in the interaction. When the light scalar distributions saturate (that is, when the gravitino production can be effective), we find numerically $a \sqrt{\langle X^2 \rangle} \sim (10^{-2} - 10^{-1}) \phi_0$. The typical comoving momenta $a p$ characterizing the scalar distributions are instead only about one order of magnitude higher than the “inflaton mass” at the end of inflation, $a p \sim 15 \sqrt{\lambda} \phi_0$. As a consequence, the gauginos acquire a mass $m_{\tilde{z}} \sim (10^2 - 10^3) p$, which shows that these processes are kinematically forbidden.²¹

We notice that, at least for this specific kind of interactions, the system is still effectively behaving as a condensate: the number densities of the scalar distributions

²¹Processes with an additional gauge interaction and the gaugino in the propagator are allowed but strongly suppressed. See the appendix for details.

the distributions themselves. This generates a high effective mass for all the particles “strongly” coupled to these scalar fields. Compare this situation with a medium in thermodynamical equilibrium: in this case both the typical momenta and the effective masses are set by the only energy scale present, namely the temperature of the system. As should be clear by the above discussion (see also [35]), the thermalization of the distributions produced at rescattering necessarily proceeds through particle fusion. Only after a sufficiently prolonged stage of thermalization, will the system be sufficiently close to thermodynamical equilibrium so to render processes as the one discussed above kinematically allowed.

As we discuss in the appendix, kinematically allowed $2 \rightarrow 2$ interactions can be obtained by taking a trilinear interaction coming from the superpotential term (27), also responsible for the Dirac mass term for the neutrinos. This can lead to processes of the kind $\tilde{N}_R X \rightarrow x \psi_{3/2}$, or $XX \rightarrow N_R \psi_{3/2}$. The physical number density of (transverse) gravitinos produced by these processes can be estimated as

$$N_{3/2}(\eta) \sim \frac{h^2}{M_P^2} \left[\frac{N_{c,X} N_{c,N}}{a^6} \right] \bigg|_{\eta=\eta_*} \left[\frac{a(\eta_*)}{a(\eta)} \right]^3, \quad \eta > \eta_*, \quad (31)$$

where η_* is the time (to be determined numerically) at which the expression in the first parenthesis reaches its maximum, while the second parenthesis is a dilution factor due to the expansion of the Universe at later times (see the appendix for details).²²

As we have remarked, this result is subject to the limit (30), where $Y_{3/2} \equiv N_{3/2}/s$, and s is the entropy density of the Universe, computed once the dominating thermal bath is formed. For practical use, we find that a more “convenient” bound can be obtained if (30) is combined with the result for the baryon asymmetry, eqn. (24). For this purpose, consider the ratio

$$\zeta \equiv \frac{Y_{3/2}}{Y_B} = \frac{23}{8\epsilon_1} \frac{N_{3/2}(\eta)}{N_N(\eta)} = \frac{23}{8\epsilon_1} \frac{N_{3/2}(\eta_*)}{N_N(\eta_*)}. \quad (32)$$

The quantity ζ has two main advantages, (i) it is independent of the entropy of the Universe and (ii) it can be computed already at $\eta = \eta_*$, since after this time the two physical number densities $N_{3/2}$ and N_N simply rescale as a^{-3} . While $Y_{3/2}$ must satisfy the upper bound (30), the limit $Y_B \gtrsim 10^{-11}$ poses a lower bound on the number density of the sneutrinos, if leptogenesis is assumed to be responsible for the baryon asymmetry of the Universe. Adopting the parameterization (4), we then see that the ratio ζ has to satisfy

$$\zeta \simeq \frac{3 \times 10^6}{\delta_{\text{CP}}} \left(\frac{10^{10} \text{ GeV}}{M} \right) \left(\frac{0.05 \text{ eV}}{m_3} \right) \frac{N_{3/2}(\eta_*)}{N_N(\eta_*)} < 10^{-3}. \quad (33)$$

It is important to stress that, unlike the limit (30), this bound cannot be ameliorated by an eventual entropy release which may occur between the decay of the r.h. sneutrinos

²²The time η_* roughly corresponds at the moment at which the distribution of light quanta X starts to saturate. This typically occurs towards the end of the rescattering stage, which guarantees that considering only $2 \rightarrow 2$ processes should provide at least an order of magnitude estimate of the gravitino quanta produced at this stage.

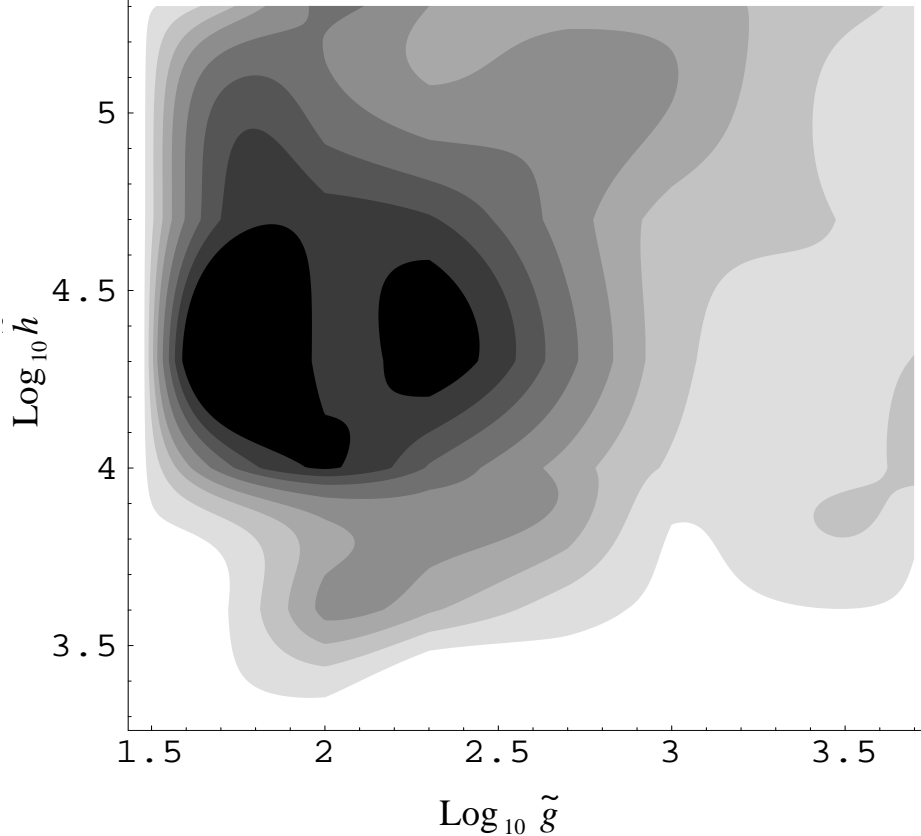


Figure 6: Contour plot for the quantity $\tilde{\zeta}/\tilde{h}$, where $\tilde{\zeta}$ is defined in eqn. (34). The contour lines range from $\tilde{\zeta}/\tilde{h} = 3 \cdot 10^{-12}$ to $\tilde{\zeta}/\tilde{h} = 1.1 \cdot 10^{-10}$ (darker region).

and nucleosynthesis, since both the gravitino and the baryon number densities would be diluted in the same amount. For this reason, we find in the present context the bound (33) more significant than the limit (30) involving the gravitino abundance alone.

We wanted to verify whether the condition (33) is respected for the choice $M = 10^{11}$ GeV considered in the previous sections, and for several values of the couplings g and h (defined in the potential (29)) in the range $\tilde{g} \in [30, 5000]$, $\tilde{h} \in [2000, 200000]$ (we remind that $\tilde{g} \equiv g^2/\lambda$, and analogously for h). To do so, we have defined

$$\zeta \equiv \frac{1}{\delta_{\text{CP}}} \left(\frac{0.05 \text{ eV}}{m_3} \right) \tilde{\zeta}, \quad (34)$$

and in figure 6 we have plotted the quantity $\tilde{\zeta}/\tilde{h}$. In this way, we factor out the explicit dependence of ζ on \tilde{h} coming from the cross section of the dominant $2 \rightarrow 2$ processes, see e.g. eqn. (31).

The qualitative behavior of $\tilde{\zeta}/\tilde{h}$ with h (vertical axis) is easily understood in terms of the arguments used to explain the results of figure 3. For relatively low h , the amplification of the X field is weak, and so few gravitinos quanta are produced. As h increases, the amplification becomes stronger, both regarding the final value at which

the number density of produced gravitinos also increases. As h further increases, the rapidity at which X saturates keeps increasing (and so, the time η_* at which most of the gravitino quanta are produced decreases); however, the final value of N_X starts to become smaller, leading to the decrease of $\tilde{\zeta}/\tilde{h}$ that we observe in the figure. For fixed values of h , the final result is also first increasing and then decreasing with g . This behavior is presumably due to the dependence on g of the total number of quanta of both the N and X fields produced at preheating/rescattering (notice that it vanishes both at very small and very high g). However, the interpretation is in this case less clear.

Concerning the value of $\tilde{\zeta}$ itself, in the range of coupling considered it reaches the maximal value at $\tilde{g} \simeq 100$, $\tilde{h} \simeq 200000$, where it evaluates to $\tilde{\zeta} \simeq 10^{-5}$. From what we have just said, we expect that $\tilde{\zeta}$ starts decreasing at higher h , although the numerical simulations we have performed show that the decrease starts only at the highest value of h that we have considered. From the definition (34), we see that the bound $\zeta < 10^{-3}$ is respected, provided the CP violation encoded in the parameter δ_{CP} (see eqs. (4) and (5)) is not too small. However, we remind that our estimates take into account only the gravitino quanta produced from the end of the rescattering stage on, while a higher production at earlier times cannot be excluded.

Acknowledgments

The numerical results presented in this paper have been obtained using the code “LATTICEASY” [36], by G. Felder and I. Tkachev. We acknowledge fruitful discussions with Patrick B. Greene, Lev Kofman, Gianpiero Mangano, and Michael Plümacher. L.B. thanks the CERN-TH division where part of this work was done. M.P. thanks the Institute for Particle Physics Phenomenology at the University of Durham for their friendly hospitality during the initial stages of this work. The work of L.B., M.P. and L.S. is supported by the European Community’s Human Potential Programme under contract HPRN-CT-2000-00152 Supersymmetry and the Early Universe. The work of L.B. was also supported in part by the Italian INFN under the project “Fisica Astroparticellare”. M.P. also acknowledges support under contracts HPRN-CT-2000-00131 Quantum Spacetime and HPRN-CT-2000-00148 Physics Across the Present Energy Frontier.

Appendix

In this appendix we derive eqn. (31) of the main text. We first estimate the number density of gravitinos produced by the nonthermal distributions formed at rescattering. We remark that at this stage supersymmetry breaking is controlled by the energy

which is “eaten” in the unitary gauge) is thus provided by a linear combination of the fermionic superpartners of these scalars, and does not coincide with the longitudinal gravitino at late times (at least for the standard case of a present gravitationally mediated supersymmetry breakdown). For this reason, our discussion will be limited to the transverse gravitino component $\psi_{3/2}$. Its mass is given by [70]

$$m_{3/2} = e^{\mathcal{K}(\phi_i)/M_P^2} \frac{|W(\phi_i)|}{M_P^2}, \quad (\text{A-1})$$

where, we remind, W and \mathcal{K} are, respectively, the superpotential and the Kähler potential of the model, while $\{\phi_i\}$ denotes the set of scalar fields amplified during preheating and rescattering. To quantize the transverse gravitino, we define a homogeneous mass $m_{3/2}^0$ by replacing in eqn. (A-1) the (\mathbf{x} -dependent) values of the fields with their (homogeneous) variances, $\phi(t, \mathbf{x}) \rightarrow \sqrt{\langle \phi^2 \rangle(t)}$. The difference $\delta m_{3/2} \equiv m_{3/2} - m_{3/2}^0$ will be accounted for in the interaction lagrangian. The main production of gravitinos is expected to occur close to the point at which the number density of light scalars X reaches its maximum, in the same way as most of the thermal production occurs as soon as the thermal bath is generated. From the results of section 4, we observe that the maximum of $N_{c,X}$ is achieved at the end of the rescattering phase. At this moment the thermalization stage begins, during which the variances of the fields show an adiabatic evolution. This allows us a consistent quantization of the gravitino component, since the mass $m_{3/2}^0(t)$ is also varying only adiabatically in this period.

It is clear that the main concern with the procedure just described is that, contrary to the thermal case, the difference $\delta m_{3/2}$ is now of the same order of $m_{3/2}^0$ itself, at least during the initial part of the thermalization stage. This leads to the problem discussed in section 5, namely to the fact that the perturbative computation of the gravitino production is presumably not as straightforward as in the thermal case, and more complicated processes than just $2 \rightarrow 2$ interactions can be expected to be relevant. However, as we have already remarked, the latter should provide at least an order of magnitude estimate for the gravitino produced from the end of the rescattering stage on, and should reasonably lead to a lower bound to the total production. For this reason, we now proceed to a rough estimate of their cross sections.

The dominant processes with two gravitinos as outgoing particles have two gravitationally suppressed vertices (i.e. $X X \rightarrow \psi_{3/2} \psi_{3/2}$ with a flat direction fermion x in the propagator; processes coming from the interaction term $\delta m_{3/2} \bar{\psi}_{3/2} \psi_{3/2}$ are subdominant). Their cross section is of the order $\sigma \sim p^2/M_P^4$, where here and in the following p denotes the typical momentum exchanged in the scattering. As in the thermal case, the distributions of the light quanta are indeed characterized by a typical momentum; while for a thermal distribution $p \sim T$, we now have $p \sim \kappa \sqrt{\lambda} \phi_0/a(t)$, where in the cases shown below κ is a coefficient of order 10 dependent on the specific choice of the parameters. In our estimates we will take $\kappa \sim 15$.²³ Thus, $\sigma \sim 10^{-11} M_P^{-2} a^{-2}$ for this class of processes.

²³The existence of a typical momentum allows the use of the integrated Boltzmann equation to estimate the amount of gravitinos produced. Moreover, since this momentum is much higher than the gravitino mass, the value of the latter does not affect the cross sections for the processes with

tionally suppressed vertex, and hence with only one gravitino in the final state. For example, the standard thermal production is mainly due to channels having a gauge interaction as the second vertex, e.g. $HH \rightarrow \psi_{3/2} \tilde{z}$ with an exchanged higgsino. In the present context, however, processes with outgoing gauginos (that we generically denote with \tilde{g}) are expected to be kinematically forbidden, since these particles acquire a high effective mass from their interaction with the nonthermal scalar distributions. Indeed, if a scalar field X has a large vev, and an interaction of the form $\sqrt{\alpha}X\psi\tilde{g}$ is present (ψ and \tilde{g} are two component matter fermion and gaugino, we use $\sqrt{\alpha}$ because we have already used g as the inflaton-neutrino coupling) then the gaugino acquires a large Dirac mass $\sim \sqrt{\alpha}\langle X \rangle$ mixing with ψ . We have large variances, rather than a large vev; by analogy with finite temperature, we expect that $\langle X^2 \rangle \neq 0$ will generate an effective mass square “ $m^2 \sim \alpha\langle X^2 \rangle$ ” in the \tilde{g} and ψ dispersion relations.²⁴ So for kinematic purposes, we assume that gauginos which couple to the flat direction have masses of order $\sqrt{\alpha\langle X^2 \rangle}$.

When the light scalar distributions saturate (that is, when the gravitino production we are discussing can be effective), we find numerically $\sqrt{\langle X^2 \rangle} \sim (10^{-2} - 10^{-1}) \phi_0/a$. As a consequence, $m_{\tilde{g}} \sim (10^2 - 10^3)p$, and these scatterings are forbidden. One is immediately led to consider processes with an additional $X_i \psi_j \tilde{g}$ vertex and in which the heavy gaugino is off-shell. Their cross section can be roughly estimated as $\sigma \sim 10^{-2} (\alpha/M_P)^2 (p/m_{\tilde{g}})^2$, which is comparable or smaller than the cross section for the process $XX \rightarrow \psi_{3/2} \psi_{3/2}$ considered above.

Finally, there is the possibility that the second vertex comes from the superpotential term (27), also responsible for the Dirac mass term for the neutrinos. This can lead to processes of the kind $\tilde{N}_R X \rightarrow x \psi_{3/2}$ or $XX \rightarrow N_R \psi_{3/2}$ (x denoting the fermionic partner of X ; all processes have in the propagator the fermionic partner of one of the incoming scalars). The cross sections for these processes are roughly estimated as²⁵ $\sigma \sim h^2/M_P^2 \simeq 10^{-13} \tilde{h}/M_P^2$.

Thus, unless of a very small coupling h , the last class of scatterings has the highest cross section and dominates the production of the transverse gravitinos. In particular, processes with one incoming N_R quantum are dominant if $N_{c,X}$ starts to saturate at a smaller value than $N_{c,N}$. Viceversa, scatterings of the kind $XX \rightarrow N_R \psi_{3/2}$ will dominate. Numerical results indicate that the former situation is more often realized. The cases in which the opposite was found are characterized by a relatively high coupling \tilde{h} , so that the light degrees of freedom are quickly amplified to values $N_{c,X} \gtrsim N_{c,N}$. Anyhow, in these cases the cross sections of the two type of processes are clearly of the same order of magnitude. Hence, for brevity of exposition we will only refer to the processes with an incoming r.h. sneutrino, although both the two

outgoing transverse gravitinos. This is welcome, since the above (somewhat arbitrary) redefinition $m_{3/2} = m_{3/2}^0 + \delta m_{3/2}$ will not affect our estimates.

²⁴It is implicit in this discussion that all the fermionic fields are quantized in the same way as done for the gravitinos.

²⁵In this estimate it is assumed that the exchanged momentum p is higher or at most comparable with the mass of the r.h. neutrinos. This is certainly true when most of the gravitinos are produced, i.e. when the distributions of light quanta X are about to saturate.

The integrated Boltzmann equation reads

$$\frac{dN_{3/2}}{dt} + 3 H N_{3/2} \simeq \langle \sigma | v | \rangle N_X N_N , \quad (\text{A-2})$$

where the “friction term” due to the expansion of the Universe can be neglected in the estimate of the order of magnitude of gravitinos produced. The whole production time can be then divided in a series of time intervals of duration $H^{-1}(t_i)$ each. During each interval, quanta of gravitinos are generated with a density of

$$\delta N_{3/2}^i \sim \frac{h^2}{M_P^2} \frac{N_{c,X}(\eta_i)}{a(\eta_i)^3} \frac{N_{c,N}(\eta_i)}{a(\eta_i)^3} H^{-1}(\eta_i) \quad (\text{A-3})$$

(notice the presence of the scale factor, since the physical and not the comoving occupation number has to be used in the integrated Boltzmann equation). The function $(N_{c,X} N_{c,N} H^{-1} a^{-6})(\eta)$ amounts to zero at the end of inflation, and it reaches a maximum at a time η_* , which can be determined numerically and which roughly corresponds to the moment at which the comoving number density $N_{c,X}$ starts saturating (this in turns occurs towards the end of the rescattering stage, when $2 \rightarrow 2$ processes start dominating). At $\eta > \eta_*$ it then quickly decreases due to the expansion of the Universe. It thus turns out that, as for the thermal production, the gravitino quanta are mostly generated at the time η_* , so that their “late time” physical number density is approximatively given by

$$N_{3/2}(\eta) \sim \frac{h^2}{M_P^2} \frac{N_{c,X}(\eta_*)}{a(\eta_*)^3} \frac{N_{c,N}(\eta_*)}{a(\eta_*)^3} H^{-1}(\eta_*) \left[\frac{a(\eta_*)}{a(\eta)} \right]^3 , \quad \eta > \eta_* . \quad (\text{A-4})$$

This is eqn. (31) of the main text.

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